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Mirror matter and heavy singlet neutrino oscillations in the early universe

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Abstract

We investigate the mixing of heavy gauge singlet neutrinos in a mirror matter model employing the seesaw mechanism. The parameter constraints that must be satisfied to prevent the overproduction of mirror matter in the early universe are deduced. We find that no fine tuning in the heavy neutral fermion sector is required for this mirror matter model to satisfy cosmological constraints. Baryogenesis scenarios are briefly discussed in the context of the mirror model.

I. INTRODUCTION

The idea that every ordinary particle is related to a mirror partner has been discussed in Refs. [1–8]. In the Exact Parity Model (EPM) of Ref. [3], mirror matter is postulated as means by which to retain invariance under all Improper Lorentz transformations, whereby each ordinary particle and its mirror partner are related via a non-standard parity symmetry. In this symmetric mirror-matter model, the parity symmetry dictates that the particle interactions in the mirror world are of precisely the same strength as those in the ordinary world. On the other hand, the authors of Ref. [6] consider non-symmetric mirror matter models in which the parity symmetry is spontaneously broken.

In the context of the EPM the ordinary and mirror sectors interact only via mixing between the ordinary and mirror neutrinos, Higgs bosons, and neutral gauge bosons, and gravitationally. The mixing between ordinary and mirror neutrinos is of great interest for explaining the observed neutrino anomalies [3,4], both because parity symmetry is a theoretically natural way to obtain the maximal mixing indicated by the Superkamiokande (SK)

experiment ¹ [10], and because the mirror neutrinos can fulfill the need for light “sterile” neutrinos. The implications of a mirror sector have also been investigated in a wide range of astrophysical and cosmological contexts, see Refs. [7,11–16].

We are concerned here with the possible oscillations amongst heavy gauge singlet neutrinos in the early universe. These are heavy right handed neutrinos (and left handed mirror neutrinos) which are employed in seesaw models of neutrino mass to suppress the masses of the light neutrinos. The essential point is that all neutrinos in the EPM (both the ordinary light neutrinos and the heavy gauge singlet neutrinos) are maximally mixed with a mirror partner. This is potentially dangerous since large amplitude oscillations between the active and mirror neutrinos in the early universe may serve to equilibrate the mirror neutrinos, even if there was initially no mirror matter present. This of course would then be in violation of constraints on the number of light particle species in thermal equilibrium at the time of big bang nucleosynthesis (BBN).

In the case of oscillations between the three light neutrinos species and their mirror partners, it turns out that the naive expectation that the mirror neutrinos are all equilibrated for $\Delta m^2 \gtrsim 10^{-6} \text{eV}^2$ can be avoided in a natural way [16]. Given an initial baryon asymmetry, there exists a large region in of parameter space for which small angle (inter-generational) ordinary-mirror mixing can amplify the asymmetry between the number of neutrinos and anti-neutrinos to be many orders of magnitude larger than the baryon asymmetry. This creates an effective potential which suppresses the maximal mixing between the ordinary and mirror neutrinos, and hence prevents the mirror species from being brought into equilibrium.

However, we have still the heavy singlet neutrinos to worry about. We will be concerned with temperatures greatly exceeding that of the electroweak phase transition, where an initial lepton or baryon asymmetry may or may not have been present depending on the specific dynamics in the early universe (for example an asymmetry may have been created by an Affleck Dine mechanism [17]). If a large asymmetry in the light lepton sector existed, it would contribute to the effective potential [18] which suppresses the oscillations of the heavy neutrinos. However, unless we were to introduce gauge bosons that couple to the singlet neutrinos, we would expect no direct feedback effect and hence no exponential growth of lepton asymmetry via oscillations (as can occur in the light neutrino sector at low temperatures), since the potential would not depend upon the asymmetry in the number of heavy neutrinos themselves. Hence, we will examine the case where no large asymmetries are responsible for suppressing oscillations.

Note that singlet neutrinos are of great interest with regard to producing the observed baryon asymmetry - either via CP violating out-of-equilibrium decays [19] or via oscillations [20]. We shall examine the interactions of the gauge singlet neutrinos and hence determine the parameter constraints that must be satisfied to render the EPM consistent with cosmology.

¹Note that results from SK [9] now appear to disfavour the $\nu_\mu \rightarrow \nu_s$ solution to the atmospheric anomaly. This is as yet preliminary and does not alter our discussion of heavy singlet neutrino dynamics. We eagerly await further results and analysis from SK on this matter.

II. MIRROR NEUTRINOS

We begin by outlining the neutrino sector of the EPM. The most natural choice to investigate is a seesaw model as was considered in [4]. We reproduce here the relevant features.

The neutrino fields consist of ν_L and N_R , and their mirror partners ν'_R and N'_L . The full gauge symmetry of the theory is

$$\text{SU}(3)_1 \otimes \text{SU}(2)_1 \otimes \text{U}(1)_1 \otimes \text{SU}(3)_2 \otimes \text{SU}(2)_2 \otimes \text{U}(1)_2, \quad (1)$$

that is, the Standard Model gauge group squared. The ν_L belongs to an $\text{SU}(2)_1$ doublet and its mirror partner ν'_R belongs to an $\text{SU}(2)_2$ doublet,

$$f_L = (\nu_L \ e_L)^T, \quad f'_R = (\nu'_R \ e'_R)^T, \quad (2)$$

while both the N_R and N'_L are gauge singlets. The Yukawa Lagrangian which includes the most general renormalisable terms allowed by the gauge and parity symmetries is

$$\begin{aligned} \mathcal{L}_{\text{Yuk}} = & \lambda_1 [\bar{f}_L \tilde{\phi}_1 N_R + \bar{f}'_R \tilde{\phi}_2 N'_L] + \lambda_2 [\bar{f}_L \tilde{\phi}_1 (N'_L)^C + \bar{f}'_R \tilde{\phi}_2 (N_R)^C] \\ & + M_1 [\bar{N}_R (N_R)^C + \bar{N}'_L (N'_L)^C] + M_2 \bar{N}_R N'_L + \text{H.c.} \end{aligned} \quad (3)$$

The ordinary $\text{SU}(2)_1$ doublet Higgs field is $\phi_1 = (\phi^+ \ \phi^0)^T$, and we employ the notation $\tilde{\phi} = i\sigma_2 \phi^*$, where σ_2 denotes a Pauli matrix. The ϕ_2 Higgs field, the parity partner of ϕ_1 , is correspondingly a doublet under $\text{SU}(2)_2$.

Below the temperature at which the electroweak phase transition occurs, the λ_1 terms produce Dirac masses for the ν 's and ν' 's, while the λ_2 mass terms lead to mixing between the ordinary and mirror neutrinos. We shall denote these masses m_1 and m_2 respectively, where $m_{1,2} = \lambda_{1,2}v$, with $v \simeq 246\text{GeV}$ being the Higgs vacuum expectation value (VEV). However, we shall be mainly interested in the symmetric phase, where the VEV vanishes and light neutrinos are massless. Note that in this case there is no mass mixing between the heavy and light sector. The M_1 and M_2 terms are bare masses, which we shall assume are much greater than the electroweak scale masses. The M_2 terms mix ordinary and mirror matter.

We shall assume that intergenerational mixing is small and consider a single generation model [3]. In the parity diagonal basis, defined by the states $(\nu_L^+ \ \nu_L^- \ (N_R^-)^C \ (N_R^+)^C)^T$ where,

$$\nu_L^\pm = \frac{1}{\sqrt{2}}(\nu_L \pm (\nu'_R)^C), \quad N_R^\pm = \frac{1}{\sqrt{2}}(N_R \pm (N'_L)^C), \quad (4)$$

the mass matrix takes the form

$$\mathcal{M} = \begin{pmatrix} 0 & 0 & 0 & m_+ \\ 0 & 0 & m_- & 0 \\ 0 & m_- & M_- & 0 \\ m_+ & 0 & 0 & M_+ \end{pmatrix}, \quad (5)$$

where $M_\pm = M_1 \pm M_2$ and $m_\pm = m_1 \pm m_2$. The light neutrino mass (and also parity) eigenstates are approximately ν_L^\pm with seesaw suppressed Majorana masses given by

$$m_a = m_+^2/M_+, \quad m_b = m_-^2/M_-, \quad (6)$$

while the heavy eigenstates N_R^\pm have masses given by

$$M_+, \quad M_-. \quad (7)$$

It is the existence of the λ_2 terms, no matter how small, that forces the mixing between an ordinary neutrino and its mirror partner to be non-zero, and hence maximal because the mass and parity eigenstates must coincide. Likewise, any nonzero value of M_2 requires that the heavy singlet neutrinos are maximally mixed.

So we see that the ordinary light neutrinos are coupled to ordinary singlets N_R via λ_1 and to mirror singlets N'_L via λ_2 . It would seem natural and, for the heavier generations, it shall turn out to be necessary, that λ_2 is smaller than λ_1 . Note that for $\lambda_1 \gg \lambda_2$, $m_\pm \simeq m_1$. We shall also assume that $M_2 < M_1$ so that M^+ and M^- have masses of approximately the same order of magnitude, which we denote as M . Hence the idea is that for each generation, we have a pair of almost degenerate mass eigenstates. Considering only one generation, the light and heavy sector neutrino masses are related according to ²

$$m_\nu \simeq \frac{m_1^2}{M} = \frac{(\lambda_1 v)^2}{M}. \quad (8)$$

III. THE EARLY UNIVERSE

Firstly, we shall be interested in the interactions which equilibrate the ordinary right-handed (and possibly also the mirror left-handed) singlet neutrinos. Let us assume that at high temperatures we have no singlet neutrinos, N_R and N'_L , in the cosmological plasma. They will then be generated by scattering processes, according to the Yukawa couplings to the light neutrinos and Higgs bosons. If the rates for these processes exceed the expansion rate $H \simeq [T(\text{GeV})]^2/10^{18}$, the N 's will achieve thermal equilibrium. We shall be working in the regime where the temperature is much larger than the mass of the Higgs bosons, so that the rates are approximately independent of these parameters.

The N 's will decay producing both ordinary and mirror light neutrinos and Higgs bosons via processes such as

$$\begin{aligned} N &\rightarrow \nu + \phi_1, \\ N &\rightarrow \nu' + \phi_2, \end{aligned} \quad (9)$$

with a total decay rate given by

$$\begin{aligned} \Gamma_D = \Gamma_{D_1} + \Gamma_{D_2} &\simeq \frac{(\lambda_1^2 + \lambda_2^2)}{16\pi} M, \quad T \lesssim M, \\ &\simeq \frac{(\lambda_1^2 + \lambda_2^2)}{16\pi} M \frac{M}{T}, \quad T \gg M. \end{aligned} \quad (10)$$

² For a three generation model, the parameter m_ν coincides with the light neutrino mass for a given generation only if we assume small intergenerational mixing.

The scattering rate for processes such as $\nu q \rightarrow N q$ (via Higgs exchange), for $T \gg M$ is given by [21]

$$\Gamma_i = \frac{9}{64\pi^3} \lambda_t^2 \lambda_i^2 T, \quad (11)$$

where $i = 1, 2$ for ordinary and mirror singlets respectively, and λ_t is the Yukawa coupling constant for the top quark. (Processes involving the top quark dominate the scattering rate since the Yukawa coupling to the Higgs boson is much larger than for the other quarks and leptons.) In addition to scattering, singlet neutrinos may be produced by inverse decay processes, with rates given by eq.(9), which however are small compared to scattering. For $T \ll M$, production of heavy neutrinos is kinematically suppressed.

The mirror singlets must not be equilibrated via scattering, as this would lead to the whole mirror sector being equilibrated, so we require $\Gamma_2 < H$ ³. This condition is most restrictive for $T = M$, leading to

$$\frac{\lambda_2^2}{M(\text{GeV})} \lesssim 2 \times 10^{-16}, \quad (12)$$

or, using eq.(8)

$$\frac{\lambda_2}{\lambda_1} < \frac{0.1}{\sqrt{m_\nu(\text{eV})}}. \quad (13)$$

We may determine if the ordinary singlets are thermalised, in terms of the size of the light masses. If the N_R 's are thermalised, we have that

$$\frac{\lambda_1^2}{M(\text{GeV})} \gtrsim 2 \times 10^{-16}, \quad (14)$$

and hence

$$m_\nu \gtrsim 10^{-2} \text{eV}. \quad (15)$$

If a light neutrino mass satisfies this bound then the corresponding heavy singlet N_R will have been thermalised. (Of course, once we go to a three generation model we need to worry about how the bases in which m and M are diagonalised are related). The atmospheric neutrino anomaly with the Superkamiokande determined mass squared difference of $\delta m^2 \sim 10^{-2} - 10^{-3}$ suggests the ν_μ mass would satisfy eq.(15). We could reasonably expect the ν_τ to be yet heavier, and in fact the EPM requires the ν_τ mass to be $m_{\nu_\tau} \sim \text{eV}$ [16] if the atmospheric problem is solved by $\nu_\mu \rightarrow \nu'_\mu$ oscillations. So we know that at least one of the ordinary singlet neutrino flavours was brought into thermal equilibrium at some stage in the early universe. We must then make sure that the amount of mirror matter produced, either by decay or oscillation of the ordinary singlets is sufficiently small.

As the dynamics of the singlet neutrinos with Yukawa coupling constants large enough that they were thermalised is qualitatively different to those with smaller Yukawa coupling constants, we shall look at these two cases separately.

³Note that when the N 's are produced, an effective potential suppresses mixing so that the flavour eigenstates and the matter-affected mass eigenstates almost coincide.

A. Thermalised singlet neutrinos

For temperatures $T \gtrsim M$, the scattering rate is greater than the decay rate, so that enforcing, as one must do, eq.(13) automatically ensures that singlet neutrino decays do not over produce mirror particles. For light neutrino masses which satisfy eq.(15), the corresponding heavy singlet decay rate will become equal to the expansion rate while the N 's are still relativistic. Hence the singlets will decay away before becoming highly non-relativistic, and the decay will not lead to appreciable reheating.

While the singlets will have all decayed before they become very non-relativistic, we shall nonetheless be concerned with oscillations that occur when the ultra-relativistic limit may not hold. For this reason, we shall give expressions for both the extreme relativistic and non-relativistic limits, and ensure both sets of bounds may safely be satisfied.

We wish to constrain the parameters of the model such that the maximal ordinary-mirror singlet oscillations are suppressed. Essentially, we need for the effective potential to be sufficiently larger than the energy difference between the vacuum mass eigenstates. The effective potential for ordinary singlets in a background medium which contains only ordinary matter, for temperatures much larger than the mass of the Higgs boson, is given by [22]⁴,

$$\begin{aligned} V(\lambda_1) &\simeq \frac{1}{8}\lambda_1^2 T, & T \gg M \\ &\simeq \frac{1}{8}\lambda_1^2 T \frac{T}{M}, & T \ll M \end{aligned} \quad (16)$$

while the potential for mirror singlets (in an ordinary matter background) is obtained by replacing λ_1 with λ_2 . It is the difference in effective potentials that is related to the oscillation parameters,

$$V = V(\lambda_1) - V(\lambda_2) \simeq V(\lambda_1). \quad (17)$$

To effectively suppress oscillation, this must be bigger than the energy difference between the two mass eigenstates, $\omega_1 - \omega_2 \simeq \Delta M^2/(2p)$ (relativistic) or $\omega_1 - \omega_2 \simeq \Delta M$ (nonrelativistic). To achieve this we will clearly require a sufficiently large λ_1 and/or a sufficiently small mass difference.

We shall determine if the necessary constraints are consistent with the light neutrino masses and mass-squared differences applicable to the resolution of the neutrino anomalies. Note that we make the approximation of a mono-energetic rather than a thermally distributed spectrum of neutrino energies, which means for $T \gg M$ we take the neutrino momentum to be given by its thermal average $p \sim T$.

⁴Note that the calculation of the finite temperature effective masses in [22] is performed in the high temperature regime $T \gg E, p$. For the regime we are interested in, namely $E \sim T$ there may be an additional coefficient of order unity in the potential (16), which however, should not greatly alter our result. See Ref. [23] for a discussion of the fermion dispersion relations without this approximation.

The rate at which mirror singlets are produced is given approximately by [24,16]

$$\Gamma(N \rightarrow N') \simeq \frac{1}{2} D \sin^2 2\theta_m, \quad (18)$$

where

$$D = \frac{1}{2} [\Gamma_{\text{scatt}}(N_R) + \Gamma_{\text{scatt}}(N'_L)] \simeq \frac{1}{2} \Gamma_{\text{scatt}}(N_R), \quad (19)$$

and the matter mixing angle θ_m is determined by

$$\sin^2 2\theta_m = \frac{\sin^2 2\theta_0}{\sin^2 2\theta_0 + [(2p/\Delta M^2)V]^2}, \quad (20)$$

with θ_0 being the vacuum mixing angle, such that $\sin^2 2\theta_0 = 1$ for maximal mixing. In eq.(20) we have assumed that the N 's are relativistic. In the nonrelativistic limit the factor $2p/\Delta M^2$ should be replaced by $1/\Delta M$.

In eq.(18), we have assumed the oscillations are adiabatic, which holds provided

$$\gamma \equiv \frac{d\theta_m}{dt} / \frac{\Delta M^2}{2p} \ll 1. \quad (21)$$

We find that away from resonance (for the period of time in question the effective potential must be large enough for the neutrinos to be away from resonance),

$$\gamma \simeq \frac{8}{\lambda_1^2} \left(\frac{T}{10^{18} \text{GeV}} \right). \quad (22)$$

At the temperature where the N 's are first thermalised

$$\gamma(T^{\text{thermalise}}) \simeq 0.01, \quad (23)$$

and since γ decreases with time, the oscillations are always adiabatic.

To estimate the constraints that must be satisfied to prevent the mirror N 's from being overproduced, we shall use the condition

$$\Gamma(N \rightarrow N') < H. \quad (24)$$

For relativistic N 's this implies

$$\frac{\Delta M^2}{M^2} < \frac{0.1\lambda_1^2}{\sqrt{m_\nu(\text{eV})}} \left(\frac{T}{M} \right)^{5/2}, \quad (25)$$

where we have used eq.(15). Given that this will be most stringent for low temperatures we shall set $T \sim M$, giving the bound

$$\frac{\Delta M^2}{M^2} < \frac{0.1\lambda_1^2}{\sqrt{m_\nu(\text{eV})}}. \quad (26)$$

In the non-relativistic limit we would have a similar condition

$$\frac{\Delta M}{M} < \frac{0.05\lambda_1^2}{\sqrt{m_\nu(\text{eV})}} \left(\frac{T}{M}\right)^{5/2} \left(\frac{\Gamma_{\text{scatt}}^{\text{rel}}}{\Gamma_{\text{scatt}}^{\text{non-rel}}}\right)^{1/2}. \quad (27)$$

However, we need only for this to hold until the N_R 's decouple, that is, when they stop being replenished by scattering. This happens at about the time they become non-relativistic anyway, so that (T/M) in eq.(27) would be at least of order 1/10, meaning that the constraint given eq.(27) is a similar condition to eq.(26). Note that for $M_2 < M_1$, $\Delta M^2/M^2 \simeq 4M_2/M_1 \simeq 2\Delta M/M$.

The criterion given in eq.(24) is useful to provide a rough estimate as to when a species is thermalised by a certain process. More accurate results could be obtained with detailed numerical work, though for our purposes approximate expressions will suffice since we are more interested in the rough size of the bounds and whether any fine tuning of parameters would be required to satisfy them. The conditions (26,27) are not particularly restrictive, and leave plenty of scope to obtain the light neutrino mass squared differences ($\delta m_{\text{light}}^2 = [m_+^2/M_+]^2 - [m_-^2/M_-]^2$) suggested by the neutrino oscillation experiments.

So we conclude that the bounds on the couplings that mix matter and mirror matter are not very severe, and no unnatural fine tuning of parameters is required to achieve consistency with cosmological constraints. This may be compared with the light neutrino sector where, similarly, no fine tuning of parameters needs to be done, due to the mechanism of asymmetry generation [16].

Finally, we wish to comment on possible scenarios for the production of a baryon and a corresponding mirror baryon asymmetry within the context of the mirror model.

We envisage a variation on the scenario proposed in ref. [20], in which neutrino oscillations create a lepton asymmetry which is then reprocessed into a baryon asymmetry by electroweak sphaleron transitions. In ref. [20], CP violating oscillations between three heavy singlet neutrino species distribute the total lepton number (which satisfies $L^{\text{total}} = 0$) unevenly between the three species. Due to a hierarchy in the Yukawa coupling constants, one of the three singlets does not communicate this asymmetry to the usual leptons before the time when sphaleron transitions freeze out, resulting in a non-zero baryon asymmetry.

In the mirror model, we similarly wish to consider CP violating oscillations, but in this case we may obtain $L \neq 0$ and $L' \neq 0$ satisfying $L + (-L') = 0$ ⁵. The heavy neutrino asymmetries are communicated to the light leptons, and the lepton and mirror lepton asymmetries are simply reprocessed into baryon and mirror baryon asymmetries respectively, with $B' = B$. Note that although the temperature of any thermalised mirror matter must be smaller than the corresponding temperature of matter, the bound is quite weak, with $T' < 0.5T$ being sufficient to satisfy BBN constraints. Hence it is possible that the time at which mirror sphalerons freeze out is not much earlier than for the ordinary sphalerons.

⁵In the EPM, oscillations amongst ordinary and mirror neutrinos conserve $L - L'$, rather than $L + L'$, because oscillations actually interchange neutrinos with mirror antineutrinos, for example $\nu_L \leftrightarrow (\nu'^C)_L$

Having ordinary and mirror baryon numbers of the same size is a nice feature, since it implies that even though at early times ordinary matter dominated, the amount of mirror matter in the universe today is the same as the amount of ordinary matter, thereby having interesting consequences for dark matter [11,12].

A detailed discussion of the parameter space where this scenario is viable may be found in ref. [20]. Basically, somewhat small singlet neutrino masses are required, $M < 100\text{GeV}$, to avoid washout of the asymmetry through lepton number violation arising from the Majorana mass. Additionally, the Yukawa coupling constants must be large enough for the asymmetry to be communicated to the usual leptons before the spheralon transitions freeze out, $\lambda_1^2 > 10^{-14}$. These bounds are consistent with phenomenologically relevant neutrino masses.

B. The lightest singlet neutrino

We shall now discuss the case where the light neutrinos and mirror neutrinos have masses

$$m_a, m_b \lesssim 10^{-2}\text{eV}, \quad (28)$$

which is quite likely to apply for electron-flavour neutrinos. In this case the singlet neutrino parameters are not subject to the bounds in subsection III A, as both λ_1 and λ_2 are small enough that the scattering production rates for singlets are always smaller than the expansion rate. There are two possibilities, either N_e and N'_e are never populated, or the N_e states (but not N'_e) are populated by processes operating at high temperatures for which the physics involved is as yet uncertain.

If the latter occurred, we may appeal to the leptogenesis scenario of out-of-equilibrium CP violating decays [19,21] to generate the baryon asymmetry of the universe. Note that baryogenesis via out-of-equilibrium decay requires singlet neutrino masses much larger than the electroweak scale, which is a completely different region of parameter space to that required for baryogenesis via neutrino oscillation as discussed above.

For leptogenesis to be successful, the N 's must be out-of-equilibrium when they decay. However, if at the time of decay they were extremely non-relativistic, significant reheating of both ordinary and mirror particles species could result, which in addition to diluting the final value of any asymmetry generated, could reduce the temperature difference between the ordinary and the mirror sector particle species.

The N_e will decay-out-of-equilibrium, but without causing appreciable reheating, if the masses of the light neutrinos are in the range

$$10^{-6}\text{eV} \lesssim m_a, m_b \lesssim 3 \times 10^{-3}\text{eV}. \quad (29)$$

Since N_e may decay into both ordinary and mirror matter, we can obtain both L and L' asymmetries:

$$\begin{aligned} L \propto \epsilon &= \frac{\Gamma(N \rightarrow \nu\phi) - \Gamma(N \rightarrow \bar{\nu}\phi^\dagger)}{\Gamma_D^{\text{total}}} \\ L' \propto \epsilon' &= \frac{\Gamma(N \rightarrow \nu'\phi') - \Gamma(N \rightarrow \bar{\nu}'(\phi')^\dagger)}{\Gamma_D^{\text{total}}} \end{aligned} \quad (30)$$

Estimating the size of ϵ and ϵ' would require assumptions about the size of CP violating phases in the Yukawa coupling constants and the singlet neutrino mass matrix. However, since at tree level, the decay rates $\Gamma(N \rightarrow \nu\phi)$ and $\Gamma(N \rightarrow \nu'\phi')$ can be of the same order of magnitude (because $\lambda_2 \sim \lambda_1$ is permitted), it would seem plausible that L' could be as large as L .

IV. CONCLUSION

We have examined bounds on the parameters in the gauge singlet neutrino sector of the EPM to ensure mirror matter is not overproduced in the early universe, and conclude that a seesaw mechanism can be implemented in the EPM with no unnatural fine-tuning of the various masses and coupling constants required.

Heavy singlet neutrinos may allow us to account for the baryon asymmetry of the universe using a leptogenesis scenario, either via CP violating oscillations or out of equilibrium decays. In both cases, even though the total energy density of mirror matter is constrained to be smaller than the energy density of ordinary matter during the radiation dominated epoch of the universe, it is conceivable that baryon and mirror baryon asymmetries of comparable size could result.

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